

Chiral Symmetry in Nuclei

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Abstract

The impact of chiral symmetry on nuclear physics is discussed in the context of recent advances in the few-nucleon systems and of dimensional power counting. The tractability of few-nucleon calculations, illustrated by very recent solutions for $A = 2 - 6$, is shown to follow from power counting based on chiral Lagrangians. The latter predicts the suppression of N -body forces, as originally shown by Weinberg. Isospin violation in the nuclear force is similarly analyzed using the results of van Kolck, and this is shown to be consistent with results from the Nijmegen phase-shift analysis. Conventional $\rho - \omega$ and $\pi - \eta$ mixing models with on-shell mixing strength are not inconsistent with naive power counting. Meson-exchange currents calculated in chiral perturbation theory are in good agreement with experiment.

1 Introduction

My talk will try to merge two rather different areas of physics: classical (conventional) nuclear physics and modern field-theory-based particle physics. This is not an easy task. A skeptic might say that if you wanted to invent a system with little apparent dynamical basis, great complexity of internal structure (spin and isospin), and almost pathological computational difficulty, you would call it a nucleus! This view was purposefully overstated, but has elements of truth nevertheless. Nuclear dynamics was mostly phenomenological for decades. The dominance of the tensor force and the one-pion-exchange potential (OPEP) means that spin and isospin play an essential role at leading order and intuition based on simple central forces doesn't often apply. This internal structure makes computational problems exceptionally challenging, which has inhibited our ability to solve (numerically) the Schrödinger equation. The net result was that in order to learn about strong-interaction dynamics in a nucleus we had to be able to calculate, and we couldn't do the latter with much accuracy.

Recent computational advances[1] in treating few-nucleon systems may have broken this logjam. In spite of all the difficulties we have finally begun to realize the potential inherent in studying few-nucleon systems. This area of physics on

which I will focus is in my opinion the biggest success story in nuclear physics in the past decade. Beginning about ten years ago, we have made spectacular progress in solving (numerically) most of the seminal problems that were discussed decades ago as crucial to the success of the field. New terminology was coined, with “exact” or “complete” denoting calculations of observables with errors of less than 1% (in spite of the computational hardships). Problems are now being solved that were considered far out of our reach ten years ago. This work is beginning to yield dynamical information, which will undoubtedly pay dividends in the future. In order to illustrate how things have changed, almost everything that we have calculated “works”, with those few disagreements with experiment being closely examined and debated and providing considerable hope for more progress.

My purview is chiral symmetry (CS) in nuclear physics. Others much more knowledgeable than I am have talked about the particle physics aspects, including the fashionable and successful Chiral Perturbation Theory (χ PT). I hope to be able to convince you that this symmetry has a dominant influence in nuclear physics. Without the symmetry, nuclear physics would be intractable. Indeed, we can and will turn this argument around: the tractability of nuclear physics provides a strong signature for the effect of CS in nuclei. Chiral perturbation theory could turn out to be the biggest advance in nuclear physics in decades, or of very limited use. There is a huge amount of information contained in our field on the behavior of the strong interactions, and we have a great opportunity for unifying all of hadronic physics. The theoretical approach used in nuclear physics unfortunately lacks (in part) the well-defined methodology of particle physics, and the challenge will be to try to change this.

I can summarize the talk by stating that CS and dimensional power counting have an opinion about: (1) the sizes of various components of the nuclear force (both isospin-conserving and isospin-violating); (2) the relative size of three-nucleon forces(3Nf); (3) the relative size of four-nucleon forces(4Nf), ... ; (4) the relative size of relativistic corrections in (light) nuclei; (5) the relative size of nucleon (impulse approximation) and meson-exchange currents in nuclear electromagnetic and weak interactions.

2 Nuclear Physics Overview

Any discussion of the role of chiral symmetry in nuclear physics must begin with a brief discussion of three topics that will tell us how nuclear theorists do business and possibly how this should change: what we do, why we do it, and what we need to

do. For the purposes of this talk when I say “nuclear physics”, I mean “low-energy few-nucleon physics”, unless stated otherwise. Restricting myself to the traditional domain of nuclear physics (\lesssim a few hundreds of MeV) frames the problem sufficiently for my allotted time.

Potentials used in the context of the Schrödinger equation (or a generalization) are central to the organization of nuclear calculations. Dynamics, assumptions, and prejudices are contained in this quantity. The reason why potentials are used is twofold and simple: (1) nuclei are self-bound configurations of nucleons, and binding cannot be achieved in (finite-order) perturbation theory; (2) when two nucleons interact and propagate between interactions there is an infrared singularity that enhances successive iterations, and this is treated exactly by the Schrödinger equation[2]. Thus our scheme is extremely efficient and reduces the complexity of calculating an amplitude to that of defining a potential.

Although the underlying dynamics of nuclei and elementary particles is shared, the two are rather different in their scales. Simply stated, nuclei are large, squishy, and soft, while particles are small, stiff, and hard. These are words used to state that the radii of nuclei follow $R \simeq 1.2A^{\frac{1}{3}}$ fm, while particles are smaller than 1 fm. The excitation energies of nuclei are typically tens of MeV or less, while particles require hundreds of MeV, and nuclear internal momenta are on average fairly small, while in particles they can be high. We can estimate the latter, \bar{p} , using the uncertainty principle in the He isotopes[3]. Equating $\bar{p}R \sim \hbar$ and $R \sim 1.5\text{-}2.0$ fm, we obtain $\bar{p}c \sim 100\text{-}150$ MeV. For mnemonic purposes only, one can equate this to the pion mass: $\bar{p}c \sim m_\pi c^2$. This is clearly inappropriate in the chiral limit, $m_\pi \rightarrow 0$. Note that this value is about half of the Fermi momentum, $\hbar k_F c = 260$ MeV, which characterizes nuclear matter. Momentum components larger than this can play a significant role in some cases and the estimate should not be taken too literally.

Given this scale for momenta there are other scales that can be constructed. Nucleons with mass M are heavy and slow moving. The average kinetic energy of a nucleon is roughly $\frac{\bar{p}^2}{M}$ or $\sim m_\pi^2/M \sim 20$ MeV, which is fairly accurate for ^2H , ^3H , ^3He , and ^4He . Because nuclei are weakly bound systems, potential and kinetic energies are comparable in magnitude. Semirelativistic calculations[4] for these nuclei (using $\sqrt{p^2 + m^2} - m$) find corrections of $\sim 5\%$ to the kinetic energy, which are typically balanced by changes in the potential. The dominant physics is nonrelativistic.

Given these scales we can easily estimate what happens when two nucleons propagate between interactions. The Green’s function, G , schematically is $1/(\bar{p}^2/M) \sim \frac{M}{m_\pi^2}$ and becomes very large for small \bar{p} . It is worth remembering that potentials (unlike amplitudes) are not uniquely defined[2]. Rather, potentials are (nonunique) subam-

plitudes, and this leads to the “off-shell” problem of nuclear physics. Although it is possible to set criteria for how one defines V , the fact that G^{-1} is small means that rather small changes in V can be compensated by the infrared singularity in G , leading to an alternative (definition of) V , which may differ substantially.

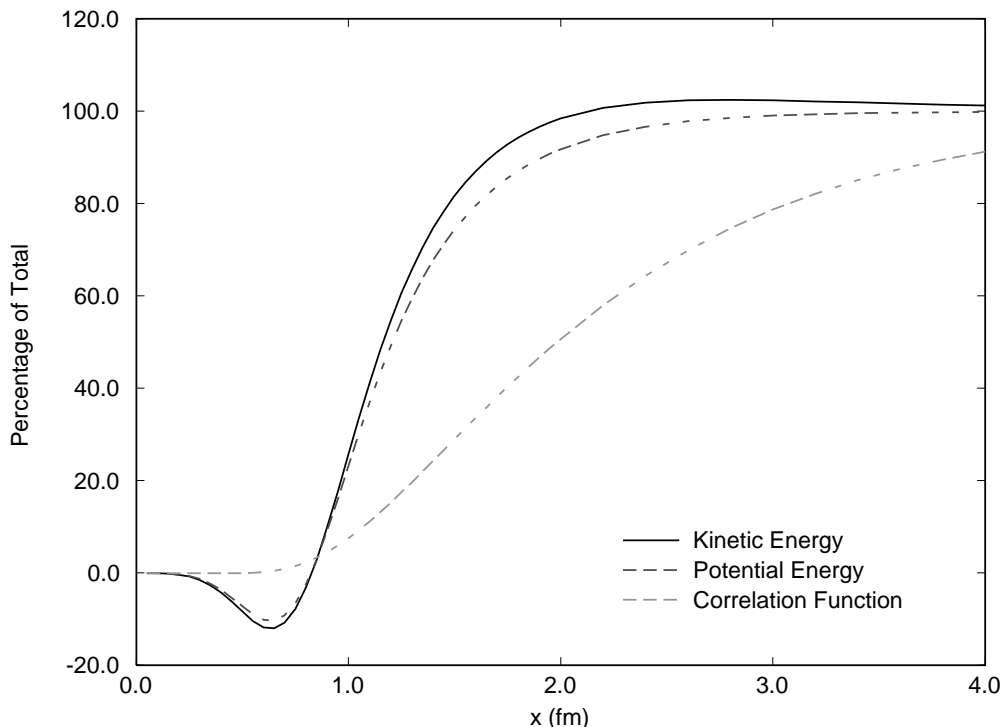


Figure 1: Percentages of accrual of kinetic energy (solid line), potential energy (short dashed line), and probability (long dashed line) within an interparticle separation, x , for any pair of nucleons in the triton.

The biggest conceptual problem in nuclear physics calculations (in my opinion) is the lack of a well-defined regularization scheme. The successive iteration of two potentials (or the sequential exchange of two mesons) is given by a loop integral, which is almost always divergent. In order to regularize this divergence nuclear potentials are cut off at short distances (large momenta), which leads to short-range repulsion and renders the calculations finite. These cutoffs are typically for momenta ~ 1 GeV, are assumed to derive from meson clouds around the nucleons (i.e., form factors), and are treated as parameters. This procedure is very efficient, physically motivated, and not very well defined, but it is the best we can do at the moment. Because of this problem very short-range operators that arise in *ab initio* calculations of potentials or transition operators are sensitive to details of the regularization procedure. Zero-range operators (which clearly require regularization) are either assumed not to contribute because of

the short-range repulsion (which makes a “hole” in the wave function), or are given a finite size. The results will be different.

Figure 1 shows the result of integrating the separation of any two nucleons in the triton out to a distance, x , and then integrating over the coordinates of the third nucleon. The three curves show the kinetic energy, potential energy, and probability (correlation function) calculated in this way, which must approach 100% as x increases. The effect of short-range repulsion and the volume element (d^3x) are obvious at small separation. Fortunately, most of the energy accrual occurs between 1 and 2 fm, corresponding to fairly small (virtual) momenta. Nevertheless, there are practical and conceptual problems at short distances in all conventional treatments. We obviously need a better way to handle this difficulty. Although the regularization question clouds the issue of testing chiral symmetry, it doesn’t change our conclusions.

3 Few-Nucleon Systems

There exists a class of recently developed nucleon-nucleon (NN) potentials that fit the available NN scattering data remarkably well and, in addition, contain much important physics[5, 6]. Many of the older potentials had a number of annoying minor defects that have been removed. These defects were highly distracting, but probably not very important in most calculations. One of these new potentials, the Argonne V_{18} [6], has 18 well-defined spin-isospin-orbital operators, illustrating the point about complexity that was made in the introduction. This potential was used to calculate the ground-state properties of the deuteron [${}^2\text{H}(1^+)$], the triton [${}^3\text{H}(\frac{1}{2}^+)$], ${}^3\text{He}(\frac{1}{2}^+)$, the α -particle [${}^4\text{He}(0^+)$], ${}^5\text{He}(\frac{3}{2}^-)$, ${}^5\text{He}(\frac{1}{2}^-)$, and ${}^6\text{Li}(1^+)$, as well as the 3^+ excited state of the latter and the (${}^6\text{He}(0^+)$, ${}^6\text{Li}(0^+)$, ${}^6\text{Be}(0^+)$) isospin triplet. Reference (7) finds the results listed in Table 1.

Table 1: Calculated and experimental ground-state energies of few-nucleon systems, together with (approximate) dates when they were first accurately solved for “realistic” potentials.

Nucleus(J^π)	${}^2\text{H}(1^+)$	${}^3\text{H}(\frac{1}{2}^+)$	${}^4\text{He}(0^+)$	${}^5\text{He}(\frac{3}{2}^-)$	${}^5\text{He}(\frac{1}{2}^-)$	${}^6\text{Li}(1^+)$
First Solved	~ 1950	1984	1987	1990	1990	1994
Expt. (MeV)	-2.22	-8.48	-28.3	-27.2	-25.8	-32.0
Theory (MeV)	-2.22	-8.47(2)	-28.3(1)	-26.5(2)	-25.7(2)	-32.4(9)

A weak three-nucleon force was added to the Hamiltonian and was adjusted to

fit the binding energy of ${}^3\text{H}$, just as the NN force fits the ${}^2\text{H}$ binding energy. The rest of the theoretical results are predictions, and are in excellent agreement with experiment. We have also indicated when the Schrödinger equation was first solved for each case. Much of the indicated progress is recent.

Examining the 2-, 3-, and 4-body cases, one finds that roughly 20 MeV/pair of nucleons accrues from the NN force, while approximately 1 MeV/triplet results from the 3Nf . If the error bar on the ${}^4\text{He}$ result is taken as an upper limit, the effect of any 4Nf should be less than .1 MeV/quartet. These numbers will be interpreted in Section 7. As spectacular as the results are, even better things are being planned. The $A = 7$ and 8 systems should be tractable when improved computers become available in the near future. These calculations demonstrate the recent achievements in the few-nucleon field.

4 Pion Degrees of Freedom

For much of its existence nuclear physics made the tacit assumption that nucleons are the only significant degrees of freedom manifested in nuclei. The reason is simple: this paradigm works, and works well. In order to demonstrate conclusively the contribution of other degrees of freedom, one must show that trustworthy calculations fail to reproduce experimental data. It was found long ago that compared to the best theoretical calculations[8] the experimental cross section for the radiative capture of thermal neutrons by protons was too large by approximately 10%. That is, calculations of this $M1$ reaction were too small if one assumed that the final photon was emitted solely from the nucleons. Moreover, since this impulse approximation is easy to evaluate, either other processes contribute or our understanding of the deuteron is seriously flawed. Since the late 1940s people had suspected that mesons were involved, but no compelling case for this scenario was made. An influential paper by Chew and Nambu[9] showed the importance of soft-pion theorems for resolving problems in our understanding of strong-interaction dynamics. Soon thereafter Riska and Brown[10] provided a compelling argument based on gauge invariance and credible phenomenology that pions were the needed ingredient for an understanding of the process. Not long thereafter (and continuing until the present time), magnetic electron scattering provided the unassailable graphic evidence for meson currents.

This is illustrated in Figs. (2) and (3)[11]. Figure (2) shows the isoscalar combination of magnetic form factors (essentially their sum) of ${}^3\text{He}$ and ${}^3\text{H}$ from a recent experimental analysis. The dashed line is an impulse approximation (nucleons only) calculation and agrees very well with the data (shaded curve). Uncharged mesons

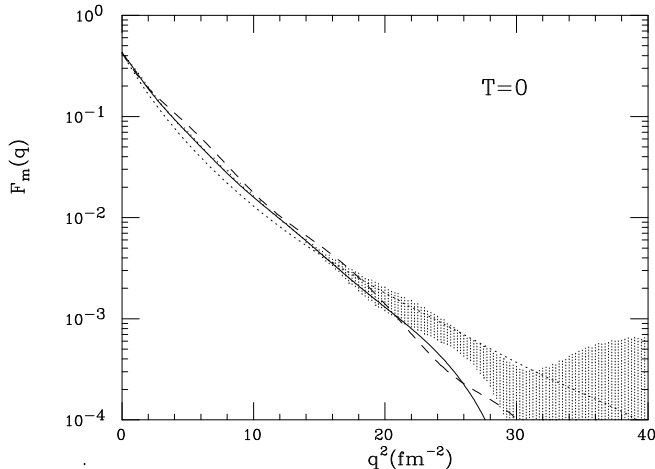


Figure 2: Isoscalar elastic magnetic form factor of the trinucleons as a function of squared momentum transfer, q^2 . The shaded area is the data, the dashed line is the impulse approximation (nucleons only), while the solid and dotted lines contain meson currents.

(corresponding to $T=0$) apparently don't play a large role. In contradistinction, the isovector combination of form factors (essentially their difference) in Fig. (3b) shows an impulse approximation calculation shifted far from the data. Inclusion of meson currents (which are dominated by single-pion exchange) corrects the problem. Figure (3a) shows the threshold (transition) magnetic form factor of the deuteron. This $^3S_1 - ^3D_1 \rightarrow ^1S_0$ reaction is just the inverse of the thermal np radiative capture. A similar pattern is found, which demonstrates conclusively the existence of pion degrees of freedom in the nucleus interacting with the external fields.

That $\Delta T = 1$ reactions should display the effect of intranuclear motion of charged mesons is perhaps no surprise, since the current continuity equation

$$\nabla \cdot \mathbf{J}_{MEC}(\mathbf{x}) = -i[V, \rho(\mathbf{x})], \quad (4.1)$$

relates the meson-exchange currents (MEC) to noncommuting parts of the potential (mostly isospin dependent). That “pure” one-pion exchange completely dominates is more surprising. This was noted by Rho[12], who called the phenomenon a “chiral filter”. Recently, a prescription was developed[13] to enforce Eq.(4.1). It can be shown that in the general case this prescription produces a nearly pointlike πNN vertex, in keeping with the “chiral filter”, and works quite well in most cases. We will discuss meson-exchange currents further in Section 7.

The final unambiguous demonstration of pion degrees of freedom in nuclei has an added cachet: it comes with error bars. Beginning approximately fifteen years ago

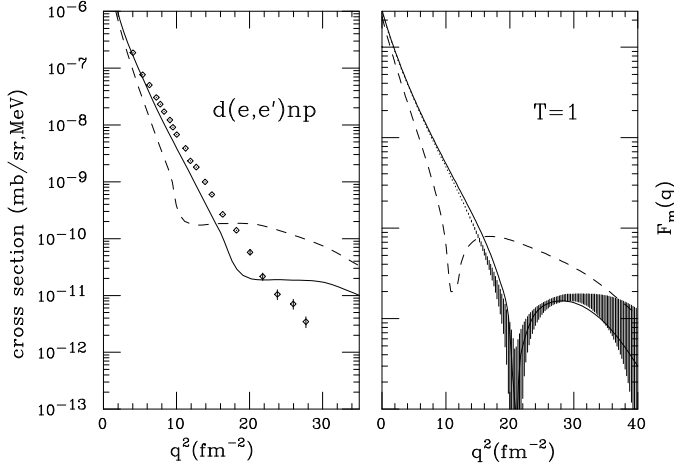


Figure 3: The (isovector) threshold deuteron electrodisintegration is shown in (a) as a function of squared momentum transfer, q^2 . The individual points are the data, the dashed line is the impulse approximation (nucleons only), while the solid line contains pion-exchange currents. The isovector elastic magnetic form factor of the trinucleons is shown in (b) as a function of squared momentum transfer, q^2 . The shaded area is the data, the dashed line is the impulse approximation (nucleons only), while the solid line contains pion-exchange currents.

the Nijmegen group[14] have implemented a sophisticated and successful program of Phase Shift Analysis (PSA) of the NN interactions. Their methodology includes treating all known long-range components of the electromagnetic interaction, such as Coulomb, magnetic moment, vacuum polarization, etc., as well as the tail of the NN interaction beyond 1.4 fm, which includes OPEP. The inner interaction region is treated in a phenomenological fashion. This allows an accurate determination of the πNN coupling constants. In order to check for systematic errors they also fit the masses of the exchanged pions, both charged and neutral, and find

$$m_{\pi^\pm} = 139.4(10) \text{ MeV} , \quad (4.2a)$$

$$m_{\pi^0} = 135.6(13) \text{ MeV} . \quad (4.2b)$$

The small error bars ($\lesssim 1\%$) demonstrate the importance of OPEP in the nuclear force. They are currently investigating the tail of the rest of the NN interaction.

A valuable byproduct of this work is the ability to construct potentials by directly fitting to the data, rather than to phase shifts, and to utilize the entire NN data base. Several potential models, such as the Argonne V_{18} model, have been constructed in this way and fit the NN data base far better than any previous attempts. One useful corollary of this work[5] is that a baseline has been set for the

triton binding energy (~ 7.62 MeV) using local NN potentials. Nonlocal potential components arising from relativity are currently under intensive investigation.

Finally, we should ask what the other consequences might be of this great sensitivity to OPE processes. In 1984 it was noted[15] that a “pure” OPEP used in certain deuteron reactions was as good as using a “realistic” potential. In the triton this force was substituted for the $^3S_1 - ^3D_1$ part of the potential and produced nearly the same binding as a realistic potential. Because that partial wave accounts for $\sim \frac{3}{4}$ of the triton potential energy, it was deduced that OPEP dominates the triton binding. This has been subsequently quantified and extended to other systems[16]. One finds that $\langle V_\pi \rangle / \langle V \rangle \sim 70\text{-}80\%$ for a wide variety of calculations ranging from “exact” treatments of the triton and α -particle to variational treatments of nuclear matter. Pion exchange is clearly of exceptional importance in nuclei, largely due to its rather long range (it is a pseudo-Goldstone boson) and spin (0^-), which produces a tensor force in leading order.

5 Nuclear and Chiral Scales and Interactions

We have already argued that the average nuclear momentum scale is $\bar{p}c \sim m_\pi c^2$. Several other scales are important. The pion mass sets the scale for chiral-symmetry breaking. The pion decay constant, $f_\pi = 92.4$ MeV, sets the scale for pion interactions. The large-mass scale, $\Lambda \sim 1$ GeV, comes from several different sources: the nucleon mass M , the masses of all heavy mesons (or resonances) such as ρ and ω , and $4\pi f_\pi$, which arises naturally in loop integrals. In any (low-energy) process constrained by chiral symmetry, we expect that the dimensionless parameter (“small momentum”/ Λ) controls the physics and, indeed, the convergence of any power-series expansion[17].

In 1990 Weinberg[18] introduced chiral perturbation theory into nuclear physics by calculating the leading-order NN force and the leading-order $3N$ force. Perhaps more important was the dimensional-power-counting scheme that he introduced and refined. The latter is a powerful tool that categorizes amplitudes (or potentials) in terms of powers of their characteristic energy or momentum scales: [“small momentum”/ Λ (large-mass scale)], as we noted above.

The leading-order nucleon-nucleon potentials are the august one-pion-exchange potential shown in Fig.(4a) and a generic short-range interaction for each spin and isospin channel (shown in Fig.(4c)), which subsumes the effect of all short-range resonant and non-resonant meson exchanges (e.g., Fig.(4b)). The leading-order three-nucleon force results from two-pion exchanges, such as those in Fig.(4d). The latter processes had been previously worked out by Coon and Friar[2] using a chiral La-

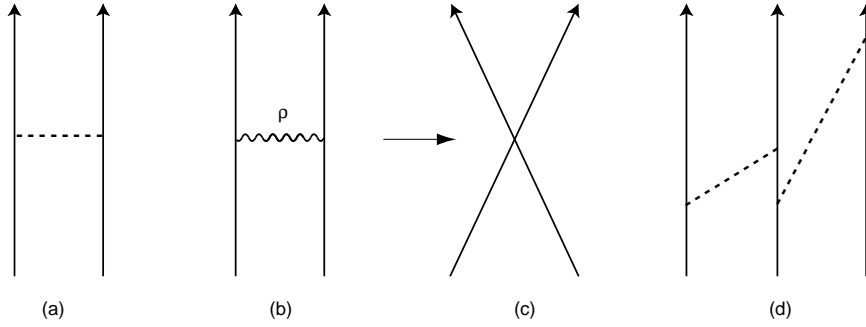


Figure 4: Time-ordered perturbation theory diagrams for nuclear potentials in χ PT, with OPEP shown in (a), ρ -exchange in (b) becomes a contact interaction in (c), while overlapping pion exchanges contribute to the 3Nf in (d). Pions are depicted by dashed lines, while nucleons are shown as solid lines.

grangian in tree order, but without considering power counting or the short-range interactions. These results appeared to differ but were later shown to be merely different off-shell extensions of OPEP[19], which is not unique, as we remarked earlier. Nevertheless, a mapping exists that exactly transforms one result into the other in leading order.

The implicit form of OPEP used by Weinberg is an energy-dependent potential, which in most many-body computational procedures is intractable. Precisely the same problem had arisen long ago in the two-nucleon problem, with competing BW (Brueckner-Watson[20]) and TMO (Taketani-Machida-Ohnuma[21]) potentials. The former, which corresponds to the Weinberg 3N force, is energy-dependent (although BW ignored those terms), while the latter is energy-independent and corresponds to the Coon-Friar 3N force. Leading-order chiral two-pion-exchange NN forces were first calculated by Ordóñez, Ray, and van Kolck[22] and subsequently verified by Friar and Coon[19]. Although this is a good start toward a nucleon-nucleon potential with proper chiral constraints, it remains to be seen whether NN scattering is sensitive to these nuances. This is being examined in Nijmegen.

Of what use is power counting? A generic Lagrangian for pions ($\vec{\pi}$), nucleons (ψ), and photons (A^μ), and containing derivatives, ∂^μ , can be pieced together with the scales we introduced earlier to produce a series of terms, each with the form

$$L \sim c_{lmn} \left[\frac{\bar{\psi}\psi}{f_\pi^2 \Lambda} \right]^l \left[\frac{\vec{\pi}}{f_\pi} \right]^m \left[\frac{\partial^\mu, A^\mu, m_\pi}{\Lambda} \right]^n f_\pi^2 \Lambda^2. \quad (5.1)$$

This can be motivated by looking at the form of the pion and nucleon masses and free energies. The dimensionless coefficients should be of order (1) if naive dimensional

power counting[23] holds, which leads to a “natural” theory. Of utmost importance is the chiral constraint[18], which is conventionally written in terms of the number of derivatives and nucleon fields at each vertex

$$\Delta = l + n - 2 \geq 0, \quad (5.2)$$

which guarantees that no Λ occurs in the numerator of Eq.(5.1). If the power series implied in Eq.(5.1) converges for nuclei and if the various c ’s are of order (1), then nuclei are “soft” and “natural”.

Several examples appropriate to nuclei illustrate these ideas. Constructing a Walecka-type model[17] with a scalar isoscalar channel and a vector isoscalar channel in zero-range form(but with no pions), one finds from Eq.(5.1)

$$L_W = \alpha_s \frac{(\bar{\psi}\psi)^2}{f_\pi^2} + \alpha_v \frac{(\bar{\psi}\gamma^\mu\psi)^2}{f_\pi^2}, \quad (5.3)$$

with values of $\alpha_s = -1.98$ and $\alpha_v = 1.48$ obtained from a recent Dirac-Hartree calculation[24]. Thus these coefficients are natural and, in fact, are quite typical.

The second example concerns the convergence of the series implied in Eq.(5.1) as a function of the nuclear density, ρ . Since $\bar{\psi}\psi \sim \rho$, and since the density of nuclear matter, $\rho_{\text{nm}} \sim 1.5 f_\pi^3$, we have

$$L \sim [c_l \sim 1][1.5 f_\pi^3 / f_\pi^2 \Lambda]^l \sim \left[\frac{1}{7}\right]^l, \quad (5.4)$$

which is fairly rapid convergence. Unfortunately, the number of terms in this expansion grows explosively as l increases. At higher densities convergence will be worse.

The third example concerns the ancient art of potential fabrication. In the old days it was found that naive PS coupling of pions and nucleons(γ_5) led to very strong forces at every order. The problem was that PS coupling optimally connects nucleon and antinucleon spinors, leading to very large and unphysical “pair” contributions to nuclear forces. An ad hoc procedure was invoked called “pair suppression”, which deleted such terms. It is easy to show[2] that these unphysical pair terms correspond to a model with $\Delta = -1$, and thus Eq.(5.2), which forbids this value, is equivalent to “pair suppression” in nuclear physics. It is also easy to show that $\Delta < 0$ leads to very strong many-body forces, which would make nuclear physics computationally intractable.

6 The Power of Counting

One typically counts powers[25] of small momenta in an amplitude: \bar{p}^ν . This procedure is quite old and is very useful, but it fails to describe the nuclear problem unless

modifications are made. A nucleus is a self-bound system that shares its available momentum; if one nucleon draws an amount from the “bank”, less is available for the others. It is this mechanism that weakens many-nucleon forces. Weinberg’s final power-counting rules[18] take this into account. A simplified version is illustrated below. Ignoring isospin factors and other factors ~ 1 , the one-pion-exchange potential has the form

$$V_\pi(\mathbf{r}) \sim \left[\frac{1}{f_\pi^2} \right] \int \frac{d^3q}{(2\pi)^3} \left[\frac{\boldsymbol{\sigma}(1) \cdot \mathbf{q} \boldsymbol{\sigma}(2) \cdot \mathbf{q}}{q^2 + m_\pi^2} \right] e^{i\mathbf{q} \cdot \mathbf{r}}. \quad (6.1)$$

The basic amplitude for OPE (in brackets) has dimension $\nu = 0$. The phase space factor has an additional dimension, $\Delta\nu = 3$. We add the two together and count this operator as having dimension $\nu = 3$. Three-nucleon operators have an additional dimension $\Delta\nu = 6$, etc. Failure to implement the momentum sharing makes it impossible to compare operators involving different numbers of nucleons for fixed A (i.e., in a given nucleus).

Finally, we can write a very simple and elegant power-counting formula for the nuclear (potential) case

$$\nu = 1 + 2(n_c + L) + \Delta, \quad (6.2)$$

where L is the number of loops, Δ is the sum of the individual Lagrangian power-counting factors ($l + n - 2 \geq 0$), and n_c is a topological factor: the number of nucleons interacting with at least one other minus the number of clusters with at least two nucleons interacting. A constant that depends only on the total number of nucleons in the nucleus has been dropped. In the most common configuration a single cluster of N nucleons interacts, which leads to $n_c = N - 1$, while in lowest order one has $L = 0$ and $\Delta = 0$. This is the leading-order N -body-force case, and we find

$$V_{Nbf} \sim \left(\frac{\bar{p}}{\Lambda} \right)^{2N-1}. \quad (6.3)$$

Thus, N -body forces weaken progressively and geometrically because of chiral symmetry ($\Delta \geq 0$). This result has enormous implications for nuclear physics.

7 Quantitative Tests and Estimates

We will examine three areas of few-nucleon physics: (1) charge-independent nuclear forces; (2) isospin-violating nuclear forces; (3) meson-exchange currents.

Dimensional power counting suggests that OPEP and the short-range nuclear forces should be comparable (both have $\nu = 3$), but they are not. The reason is simple. Strong short-range repulsion produces a coherent effect (a barrier, whose penetration is suppressed). This in turn produces a hole in the wave function which always

“wins”, no matter how strong the potential, and OPEP therefore dominates. Two-pion-exchange(TPE) is predicted to be suppressed because it corresponds to $\nu = 5$. Although most realistic potentials have a phenomenological component of two-pion range, there is not yet any direct quantitative evidence that those components play a significant role in NN scattering (note that we are not discussing ρ -meson exchange, but rather the long-range tail of the force in the TPE channel). Tests of the TPE potential are planned at Nijmegen.

The most important evidence for chiral symmetry in nuclei is that nuclear physics is tractable. The recent[7] calculation of $A = 2 - 6$ used a realistic NN force model and a weak $3N$ force adjusted to fit ${}^3\text{H}$; it also fits ${}^4\text{He}$. Their results are in good accord with power-counting predictions

$$\langle V_{NN} \rangle \sim 20 \text{ MeV/pair} , \quad (7.1a)$$

$$\langle V_{3Nf} \rangle \sim 1 \text{ MeV/triplet} , \quad (7.1b)$$

$$\langle V_{4Nf} \rangle \lesssim 0.1 \text{ MeV/quartet} , \quad (7.1c)$$

since (leading-order) NN forces correspond to $\nu = 3$, $3N$ forces to $\nu = 5$ and $4N$ forces to $\nu = 7$. The latter are therefore likely to be negligible.

Finally, we note that relativity appears to be a correction in few-nucleon systems. While not specifically a consequence of chiral symmetry, an expansion in powers of $1/M \sim 1/\Lambda$ is intimately related to the usually power counting. We note that regularizing (in the nuclear physicist’s fashion) at a momentum $\sim \Lambda$ generates relativistic corrections $\sim 5\%$ to nuclear energies. This result is somewhat controversial and requires more study (particularly using fully-relativistic calculations), but is unlikely to be qualitatively wrong.

Isospin violation in the nuclear force is still a rather poorly understood phenomenon. The isospin dependence of the nuclear force is classified according to four categories[26], with two-nucleon operators having forms

- | | |
|--|------------------------------------|
| (I). 1 and $\mathbf{t}_1 \cdot \mathbf{t}_2$ | charge independent (CI) |
| (II). $t_1^z t_2^z - \mathbf{t}_1 \cdot \mathbf{t}_2 / 3$ | charge-independence breaking (CIB) |
| (III). $(t_1 + t_2)_z$ | charge-symmetry breaking (CSB) |
| (IV). $(t_1 - t_2)_z$ and $(\mathbf{t}_1 \times \mathbf{t}_2)_z$ | charge-symmetry breaking (np only) |

where \mathbf{t}_1 and \mathbf{t}_2 are the isospin operators of nucleons “1” and “2”. The class (II) operator is an isotensor, while (III) and (IV) are isovectors. The class (IV) operators are nonvanishing only for the np system, while class (III) vanishes for the np system.

It was shown by van Kolck[25] that $I > II > III > IV$. He did this by constructing effective chiral Lagrangians corresponding to $d - u$ quark mass differences, characterized by $\epsilon = \frac{m_d - m_u}{m_d + m_u} \sim 0.3$, which has the isospin character of t_z . A separate construction of effective Lagrangians resulting from freezing out hard-photon exchanges is characterized by the fine-structure constant, α . These two Lagrangian “ladders” (each corresponding to different powers of small momenta) must be re-aligned so that the separate dimensionless factors of ϵ and α are taken into account. Using the mass shifts of neutral and charged mesons and nucleons as input, van Kolck argued that quark-mass terms were dominant in leading order. The relevant parts of his Lagrangian are

$$L \sim -\frac{g_A}{f_\pi} \bar{N} \boldsymbol{\sigma} \cdot \nabla \mathbf{t} \cdot \boldsymbol{\pi} N + [\text{mass} - \text{splitting terms}] + \frac{\beta_1}{2f_\pi} \bar{N} \boldsymbol{\sigma} \cdot \nabla \pi_0 N \\ + \gamma_s \bar{N} t_3 N \bar{N} N + \gamma_\sigma \bar{N} \boldsymbol{\sigma} t_3 N \cdot \bar{N} \boldsymbol{\sigma} N. \quad (7.2)$$

The first term is the usual πNN coupling that conserves isospin, while the remaining terms violate isospin. The very important mass-splitting interactions for charged and neutral pions and nucleons are next, followed by an isospin-violating $\pi^0 NN$ coupling and two short-range NN terms. Nuclear physicists have approached calculations of isospin-violating mechanisms in terms of resonance-saturation diagrams of the type shown in Fig.(5a). Pseudoscalar-meson exchange[27] can incur mixing (at the blob) that converts a π^0 to an η (or η'). In the limit of large η mass this produces β_1 , which should be on the order of $(\epsilon m_\pi^2/\Lambda^2)$. Numerical evaluation of the mixing diagram using typical constants leads to $\beta_1 \sim \frac{2}{3}(\epsilon m_\pi^2/\Lambda^2)$, which is a reasonable result. The same procedure can be applied to vector-meson exchange[27] and leads to the $\rho - \omega$ mixing depicted in Fig.(5b) and to γ_s , which should be of order $(\epsilon m_\pi^2/f_\pi^2 \Lambda^2)$. Evaluating γ_s in terms of the $\rho - \omega$ mixing model with typically used parameters[27, 28] one finds $\gamma_s \sim \frac{1}{2}(\epsilon m_\pi^2/f_\pi^2 \Lambda^2)$, a not unreasonable value. There have been many recent claims that the $\rho - \omega$ mixing parameter (fitted on-shell for $q^2 \sim m_\rho^2$) has a strong off-shell suppression when taken into the nuclear regime ($q^2 \leq 0$). Various calculations ranging from QCD sum rules to quark models have advanced this claim[29]. Chiral Perturbation Theory, on the other hand, makes no statements about specific models.

Finally, pseudovector-meson exchange can take place (Fig.(5c)) and leads to a spin-spin force in leading order. Close-lying isospin doublets such as (a_1, f_1) could contribute substantially to the γ_σ -term. To the best of our knowledge no such mechanism has ever been proposed, but it arises quite naturally.

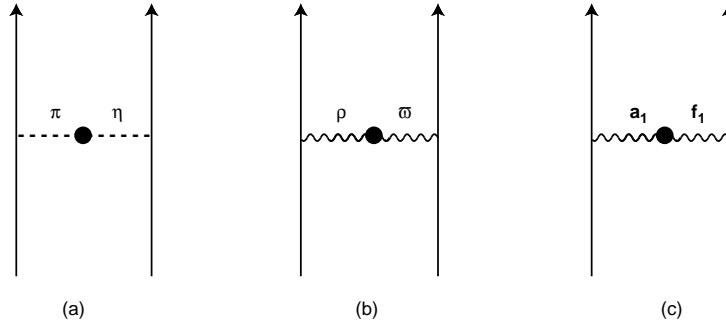


Figure 5: Meson-mixing processes that contribute to the isospin-violating NN force, including pseudoscalar-meson exchange in (a), vector-meson exchange in (b), and axial-vector-meson exchange in (c).

Given a variety of isospin-violating mechanisms (which are all $\lesssim 1\%$ of the isospin-conserving force), which of them are the largest and how important are they in nuclear physics? Because OPEP is dominant the pion-mass-difference term is very large ($\sim 3\%$ of OPEP) and leads to a large class II force. Charge-symmetry breaking of class III is next in size and is of both conceptual and practical importance in the few-nucleon problem. Class IV forces require antisymmetric complementary spin-space operators and are therefore the smallest.

Class III forces are a very important aspect of the NN interaction and are part of one of the most important success stories in few-nucleon physics. The ${}^3\text{He} - {}^3\text{H}$ mass difference can be written as the difference of the individual nucleon masses plus a (positive) binding energy difference, with ${}^3\text{He}$ being less bound than ${}^3\text{H}$. Because ${}^3\text{He}$ has a pp pair and ${}^3\text{H}$ an nn pair (each having two np pairs), we expect that most of the 764 keV binding energy difference is due to the Coulomb interaction between the two protons in ${}^3\text{He}$, which generates 648(4) keV. Small contributions from the repulsive magnetic moment interactions, the motion of the protons, the $n - p$ mass difference in the kinetic energy, and similar small mechanisms generate 45(5) keV. The remaining contribution has a short-range nature. After removal of electromagnetic effects from their interactions the $T = 1$ s-wave scattering lengths of nucleons (${}^1\text{S}_0$ state) have the values[28]: $np(-23.7 \text{ fm})$, $pp(-17.3(4) \text{ fm})$, $nn(-18.8(3) \text{ fm})$. The difference of nn and pp is $-1.5(5) \text{ fm}$, which generates approximately 66(22) keV binding energy difference[7]. The total[30] of 759(25) keV agrees well with experiment: 764 keV. This impressive success in few-nucleon physics does, however, raise one significant question. If $\rho - \omega$ mixing was greatly overestimated in the past, are there enough other mechanisms of sufficient size to account for the $-1.5(5) \text{ fm}$ difference in scattering lengths?

In recent years the Nijmegen PSA has successfully measured the πNN coupling constants, of which there are three, by focusing on the long-range part of the nuclear force[14]. Defining

$$f^2 = \left(\frac{g_A m_{\pi^+} d}{2f_\pi} \right)^2 / 4\pi, \quad (7.3)$$

where $d - 1$ is the Goldberger-Treiman discrepancy[31], one can transform measurements of f^2 into measurements of d . The exchange of a neutral pion between a pair of protons generates $f_{\pi^0 pp}^2 = 0.0751(6)$, while a neutral pion exchanged between an np pair gives $f_{\pi^0 nn} f_{\pi^0 pp} = 0.0752(8)$. Charged-pion exchange in the latter case gives $f_{\pi^+ np}^2 = 0.0741(5)$. These three pieces of experimental information can be analyzed in terms of three other quantities: the isospin-symmetric d ; the quantity β_1 ; a consistency condition, c , which should equal 1. One finds[32]

$$d - 1 = 2.0(5)\%, \quad (7.4a)$$

$$\beta_1 = 1(8) \cdot 10^{-3}, \quad (7.4b)$$

$$c = 1.007(6). \quad (7.4c)$$

The value of $d - 1$ is considerably smaller than older values and corresponds to a (monopole) form factor mass of ~ 1 GeV. Dimensional estimates of β_1 are $\sim 6 \cdot 10^{-3}$, consistent with Eq.(7.4b); Eq.(7.4c) is also satisfactory. Because a huge amount of data was analyzed to obtain these results, it will be extremely difficult to do better.

Finally, we complete the circle with a discussion of meson-exchange currents. As discussed thoroughly by Park, Min, and Rho[8], it is impossible to understand the isovector magnetic form factors in nuclei without incorporating the effect of pion degrees of freedom. The same observation holds for the time component of the nuclear axial current[33]. A number of processes contributing to nuclear electromagnetic interactions are shown in Fig.(6).

What about the effect of short-range (e.g., ρ) meson exchange? This question was answered by Rho[12], who noted that power counting for meson-current operators has the form

$$\nu = 1 + 2(n_c + L) + \Delta + \delta, \quad (7.5)$$

where $\Delta \geq -1$ for an electromagnetic vertex (but ≥ 0 for a strong vertex) and δ is a spinor-reduction factor that equals 0 (for “even” operators) or 1 (for “odd” operators). Specializing to the space part of the vector current, we find that the impulse approximation ($\sim \gamma$, corresponding to $\delta = 1$) has $\nu = 1$ in leading order and is shown in Fig.(6a). The usual seagull MEC in Fig.(6b) has $\nu = 2$, and the pion-exchange diagrams are therefore the largest corrections to the impulse approximation.

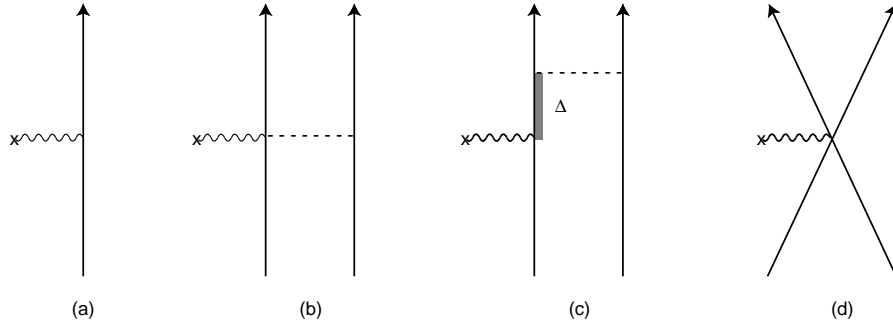


Figure 6: Nuclear electromagnetic interactions, with the impulse approximation shown in (a), seagull MEC depicted in (b), isobar-mediated MEC illustrated in (c), and short-range MEC sketched in (d). Solid lines are nucleons, dashed lines are pions, and wavy lines are (virtual) photons.

Heavy MEC (Fig.(6d)) and isobar-mediated pion-range MEC (Fig.(6c)) have $\nu = 4$. The suppression of heavy-meson exchange arises because insertion of a photon into a meson propagator involves cutting that propagator, thus making two of them. The extra propagator $\sim 1/m_\rho^2 \sim 1/\Lambda^2$, and this suppresses Fig.(6d) by two powers of ν compared to Fig.(6b). This is a lovely result!

Park, Min, and Rho[8] calculate all contributions to the (magnetic dipole) np radiative capture transition with $\nu \leq 4$ and obtain excellent agreement with experiment. In view of the complexity of the calculation and its close relationship to chiral perturbation theory calculations in the one-nucleon sector, this is an important technical achievement.

8 Conclusions

One-pion exchange dominates in the binding of light nuclei and in meson-exchange currents. This follows from power counting and the coherent suppression resulting from barrier penetration at short distances. Chiral symmetry provides order in nuclear forces: without this symmetry nuclear physics would be intractable. Turning the argument around, the tractability of nuclear physics provides strong evidence for chiral symmetry, which weakens N -body forces as N increases and n -pion exchanges compared to OPEP. The sizes of various meson-exchange currents can be understood in terms of dimensional power counting. Mechanisms for isospin violation in the nuclear force are also consistent with dimensional power counting. Finally, few-nucleon systems continue to be the testing ground for new ideas in nuclear physics because of

our ability to calculate accurately in those systems.

9 Acknowledgements

This work was performed under the auspices of the U. S. Department of Energy. Numerous discussions with J. de Swart were very helpful. I. Sick generously contributed figures. I would like to dedicate this talk to Gerry Brown, who was the first to argue that chiral symmetry was an essential ingredient of nuclear forces[34], and to Bryan Lynn[17] who has argued that CS provides a framework for nuclear physics.

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